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NO VAN DAM-VELTMAN-ZAKHAROV DISCONTINUITY FOR SUPERGRAVITY IN ADS SPACE

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Abstract

Adding explicit mass terms for the spin 2 and spin 3/2 field of $N = 1$ anti-de Sitter supergravity, the limit $M^2 \rightarrow 0$ for these mass terms is smooth: there is no van Dam-Veltman-Zakharov mass discontinuity in the propagators when the cosmological constant is non-vanishing.

Recently Kogan *et al.* [1] and Porrati [2] have shown that the mass discontinuity arising in the massless limit of massive gravity theories is peculiar to Minkowski space and does not arise in anti-de Sitter spaces. A similar result for de Sitter spaces was obtained by Higuchi [3], but in this case there are fewer physical applications (there are no unitary representations [2], and supergravity does not exist in de Sitter space [4]).

The action they consider is a sum of the Einstein action, a cosmological term Λ and the spin 2 Fierz-Pauli [5] mass term

$$\begin{aligned}\mathcal{L}_E &= \frac{\sqrt{-\det(g+h)}}{2} [R(g+h) - 2\Lambda + h_{\mu\nu}T^{\mu\nu}] + \mathcal{L}_{(mass)}^{(2)}, \\ \mathcal{L}_{(mass)}^{(2)} &= -\frac{\sqrt{-g}M^2}{8} (h_{\mu\nu}h_{\rho\sigma}g^{\mu\rho}g^{\nu\sigma} - h_{\mu\nu}h_{\rho\sigma}g^{\mu\nu}g^{\rho\sigma}).\end{aligned}\quad (1)$$

Here $g_{\mu\nu}$ is the background metric for the Einstein space. All indices are raised, lowered and contracted with the background metric $g_{\mu\nu}$, and the quantum gravitational field $h_{\mu\nu}$ is coupled to an external source $T_{\mu\nu}$ which is covariantly conserved in the background metric.

The propagator with both Λ and M non-vanishing has only a pole at $\nabla^2 = M^2 - 2\Lambda$ whose residue is [1, 2]

$$T^{\mu\nu}T_{\mu\nu} + T \left[\frac{\Lambda - M^2}{3M^2 - 2\Lambda} \right] T, \quad (2)$$

where T denotes the trace of $T^{\mu\nu}$. For $\Lambda \rightarrow 0$ at fixed M^2 , one finds $T^{\mu\nu}T_{\mu\nu} - \frac{1}{3}T^2$ which is the residue of a massive spin 2 particle. This residue is positive and thus tree level unitarity is preserved [6]. However, for $M^2 \rightarrow 0$ at fixed Λ one finds instead $T^{\mu\nu}T_{\mu\nu} - \frac{1}{2}T^2$ which is the residue of a massless spin 2 particle, and which also satisfies the tree level unitarity conditions [6]. The fact that $T^{\mu\nu}T_{\mu\nu} - \frac{1}{3}T^2$ is discreetly different from $T^{\mu\nu}T_{\mu\nu} - \frac{1}{2}T^2$ is the well-known van Dam-Veltman-Zakharov (vDVZ) mass-discontinuity [7]. As observed in [1, 2], the discontinuity is an accident in Minkowski space, but the limit $M \rightarrow 0$ is smooth in anti-de Sitter space ($\Lambda < 0$) [1, 2] and de Sitter space ($\Lambda > 0$) [3].

In this note we extend this result to spin 3/2 Rarita-Schwinger fields. Then we put the results into the context of supergravity with super-cosmological constant. Finally, we also generate a mass term through a generalized Stückelberg formalism.

The action for a massive spin 3/2 field in curved space is [8]

$$\mathcal{L}^{3/2} = -\frac{e}{2}\bar{\Psi}_\mu\gamma^{\mu\rho\sigma}\nabla_\rho\Psi_\sigma + \frac{M}{2}\bar{\Psi}_\mu\gamma^{\mu\nu}\Psi_\nu + \bar{\Psi}^\mu J_\mu, \quad (3)$$

where $\gamma^\mu = e^\mu_m\gamma^m$ with constant γ^m , $e = \det(e^\mu_m)$, $\gamma^{\mu\nu} = \frac{1}{2}[\gamma^\mu, \gamma^\nu]$ and $\gamma^{\mu\nu\rho}$ is the totally antisymmetric part of $\gamma^\mu\gamma^\nu\gamma^\rho$. Further, Ψ_μ is a Majorana spinor, so $\bar{\Psi}_\mu = \Psi_\mu^T C$ with C the charge conjugation matrix, and

$$\nabla_\rho\Psi_\sigma = \partial_\rho\Psi_\sigma + \frac{1}{4}\omega_\rho^{mn}(e)\gamma_{mn}\Psi_\sigma - \Gamma_{\rho\sigma}^\tau\Psi_\tau, \quad (4)$$

where $\Gamma_{\rho\sigma}^\tau$ is the Christoffel connection and $\omega_\rho^{mn}(e)$ the spin connection. The source J_μ is an external source, which we take to be covariantly conserved; this is necessary in the massless theory [8], but not necessary for the massive theory, although Kaluza-Klein theories automatically couple massive fields to conserved sources [1].

The field equations read

$$\gamma^{\mu\rho\sigma}\nabla_\rho\Psi_\sigma - M\gamma^{\mu\nu}\Psi_\nu = J^\mu. \quad (5)$$

Using $\gamma^{\mu\rho\sigma} = \gamma^\mu\gamma^\rho\gamma^\sigma - g^{\mu\rho}\gamma^\sigma - g^{\rho\sigma}\gamma^\mu + g^{\mu\sigma}\gamma^\rho$, we obtain

$$\gamma_\mu \not{\nabla}\gamma \cdot \Psi - \nabla_\mu \gamma \cdot \Psi - \gamma_\mu \nabla \cdot \Psi + \not{\nabla}\Psi_\mu - M\gamma_{\mu\nu}\Psi^\nu = J_\mu, \quad (6)$$

where we used that the derivative ∇_ρ in (4) commutes with γ_σ , so $[\nabla_\rho, \gamma_\sigma] = 0$. To obtain explicit expressions for the lower spin parts $\gamma \cdot \Psi$ and $\nabla \cdot \Psi$ in terms of J^μ , we contract the field equations with ∇^μ and γ^μ , respectively. This yields

$$\not{\nabla} \not{\nabla}(\gamma \cdot \Psi) - \nabla^2(\gamma \cdot \Psi) - \not{\nabla}(\nabla \cdot \Psi) + \nabla_\mu(\not{\nabla}\Psi^\mu) - M\gamma^{\mu\nu}\nabla_\mu\Psi_\nu = \nabla \cdot J = 0, \quad (7)$$

$$2 \not{\nabla}(\gamma \cdot \Psi) - 2 \nabla \cdot \Psi - 3 M \gamma \cdot \Psi = \gamma \cdot J, \quad (8)$$

where we used $\gamma^\mu \not{\nabla}\Psi_\mu = - \not{\nabla}\gamma \cdot \Psi + 2 \nabla \cdot \Psi$ in (8). In (7), we used

$$\nabla_\mu \not{\nabla}\Psi^\mu = \not{\nabla}\nabla \cdot \Psi + \gamma^\nu[\nabla_\mu, \nabla_\nu]\Psi^\mu. \quad (9)$$

We now need some properties of gravitationally covariant derivative for spinors. First of all, from the vielbein ‘‘postulate’’ $\nabla_\nu e_\rho^m \equiv \partial_\nu e_\rho^m - \Gamma_{\nu\rho}^\sigma e_\sigma^m + \omega_\nu^m{}_n e_\rho^n = 0$ we obtain $[\nabla_\mu, \nabla_\nu]e_\rho^m = 0$, which relates the curvature in term of Christoffel symbols to the curvature in term of the spin connection

$$-R_{\mu\nu\rho}{}^\sigma(\Gamma)e_\sigma^m + R_{\mu\nu}{}^m{}_n(\omega)e_\rho^n = 0, \quad (10)$$

where

$$\begin{aligned} R_{\mu\nu\rho}{}^\sigma(\Gamma) &= \partial_\mu \Gamma_{\nu\rho}^\sigma + \Gamma_{\mu\tau}^\sigma \Gamma_{\nu\rho}^\tau - (\mu \leftrightarrow \nu), \\ R_{\mu\nu}{}^m{}_n(\omega) &= \partial_\mu \omega_\nu^m{}_n + \omega_\mu^m{}_k \omega_\nu^k{}_n - (m \leftrightarrow n). \end{aligned} \quad (11)$$

We define anti-de Sitter space by

$$\begin{aligned} R_{\mu\nu\rho\sigma}(\Gamma) &= \frac{\Lambda}{3}(g_{\mu\sigma}g_{\nu\rho} - g_{\mu\rho}g_{\nu\sigma}), \\ R_{\mu\sigma}(\Gamma) &= R_{\mu\nu\rho\sigma}(\Gamma)g^{\nu\rho} = \Lambda g_{\mu\sigma}. \end{aligned} \quad (12)$$

Next we recall that, although the Christoffel symbols cancel in the curls $\nabla_\rho\Psi_\sigma - \nabla_\sigma\Psi_\rho$ in the action, they are still present in $\nabla_\mu \not{\nabla}\psi^\mu$. Hence,

$$[\nabla_\mu, \nabla_\nu]\Psi^\mu = -R_{\mu\nu}{}^{\mu\tau}(\Gamma)\Psi_\tau + \frac{1}{4}R_{\mu\nu}{}^{mn}(\omega)\gamma_{mn}\Psi^\mu = \Lambda\Psi_\nu + \frac{\Lambda}{6}\gamma_{\mu\nu}\Psi^\mu. \quad (13)$$

Similarly,

$$\begin{aligned} \not{\nabla} \not{\nabla}(\gamma \cdot \Psi) &= \nabla^2(\gamma \cdot \Psi) + \frac{1}{2}\gamma^\mu\gamma^\nu[\nabla_\mu, \nabla_\nu]\gamma \cdot \Psi = \\ &= \nabla^2(\gamma \cdot \Psi) + \frac{1}{8}\gamma^\mu\gamma^\nu R_{\mu\nu}{}^{mn}(\omega)\gamma_{mn}(\gamma \cdot \Psi) = \nabla^2(\gamma \cdot \Psi) - \Lambda(\gamma \cdot \Psi). \end{aligned} \quad (14)$$

With these results, the contracted equations reduce to

$$\frac{\Lambda}{2}\gamma \cdot \Psi + M(\not{\nabla}(\gamma \cdot \Psi) - \nabla \cdot \Psi) = 0, \quad (15)$$

$$2 \not{\nabla}(\gamma \cdot \Psi) - 2 \nabla \cdot \Psi - 3 M \gamma \cdot \Psi = \gamma \cdot J. \quad (16)$$

The first equation allows to express $\nabla \cdot \Psi$ in terms of $\gamma \cdot \Psi$, and substituting the result into (15) yields an expression for $\gamma \cdot \Psi$ in terms of $\gamma \cdot J$:

$$\nabla \cdot \Psi = \left(\frac{\Lambda}{2M} + \nabla \right) \gamma \cdot \Psi, \quad (17)$$

$$\gamma \cdot \Psi = - \left(\frac{\Lambda}{M} + 3M \right)^{-1} \gamma \cdot J. \quad (18)$$

Substituting these results into the field equation (6) leads to the propagator

$$\bar{J}^\mu \Psi_\mu = \bar{J}^\mu \left\{ \frac{(\Lambda + 3M^2)^{-1}}{\nabla + M} \left[-M \nabla_\mu \gamma_\nu - \left(\frac{\Lambda}{2} + M^2 \right) \gamma_\mu \gamma_\nu + g_{\mu\nu} (\Lambda + 3M^2) \right] \right\} J^\nu. \quad (19)$$

For fixed Λ but $M \rightarrow 0$ we obtain

$$\bar{J}^\mu \frac{1}{\nabla} \left[g_{\mu\nu} - \frac{1}{2} \gamma_\mu \gamma_\nu \right] J^\nu. \quad (20)$$

This is the propagator for a massless spin 3/2 particle. In flat space the residue can be written as $\bar{J}^\mu \left(\eta_{\mu\nu} \not{\partial} + \frac{1}{2} \gamma_\mu \not{\partial} \gamma_\nu \right) J^\nu$ which is positive definite, and the theory is thus unitary at tree level [8]. On the other hand, for fixed M but $\Lambda \rightarrow 0$, we obtain

$$\bar{J}^\mu \frac{1}{\nabla + M} \left[-\frac{1}{3M} \nabla_\mu \gamma_\nu + g_{\mu\nu} - \frac{1}{3} \gamma_\mu \gamma_\nu \right] J^\nu. \quad (21)$$

In the flat space ∇_μ commutes with $(\nabla + M)^{-1}$ and annihilates on \bar{J}^μ , and we find the propagator of a massive spin 3/2 particle in flat space [8]

$$\bar{J}^\mu \frac{1}{\square - M^2} \left[\eta_{\mu\nu} (\not{\partial} - M) + \frac{1}{3} \gamma_\mu (\not{\partial} + M) \gamma_\nu \right] J^\nu. \quad (22)$$

Also this propagator has a positive definite residue, and thus satisfies tree level unitarity [8].

The fact that $\eta_{\mu\nu} \not{\partial} + \frac{1}{2} \gamma_\mu \not{\partial} \gamma_\nu$ is discreetly different from $\eta_{\mu\nu} (\not{\partial} - M) + \frac{1}{3} \gamma_\mu (\not{\partial} + M) \gamma_\nu$ for $M \rightarrow 0$ is a van Dam-Veltman-Zakharov mass discontinuity, but in curved space with non-vanishing Λ the limit $M \rightarrow 0$ is smooth. Thus spin 3/2 fields and spin 2 fields behave the same way as far as the limit $M \rightarrow 0$ is concerned, which are reasons to turn to supergravity.

First we consider what happens if one does not take the flat-space limit. In curved space, ∇^μ does, of course, not commute with $(\nabla + M)^{-1}$. It is possible to rewrite $(\nabla + M)^{-1}$ as $[(\nabla - M)(\nabla + M)]^{-1} (\nabla - M)$. Furthermore, the operator $[(\nabla - M)(\nabla + M)]^{-1}$ is the inverse of $[(\nabla - M)(\nabla + M)]$ and the latter satisfies

$$[(\nabla - M)(\nabla + M), \nabla_\mu] \gamma \cdot J = \Lambda \nabla_\mu (\gamma \cdot J). \quad (23)$$

Acting on a vector spinor J_μ , it reads

$$\begin{aligned} [(\nabla - M)(\nabla + M)] J_\nu &= \left(\nabla^2 - M^2 \right) J_\nu + \frac{1}{2} \gamma^\rho \gamma^\sigma \left[-R_{\rho\sigma}{}^\tau{}_\tau (\Gamma) J_\tau + \frac{1}{4} R_{\rho\sigma}{}^{mn}(\omega) \gamma_{mn} J_\nu \right] \\ &= \left(\nabla^2 - M^2 - \Lambda \right) J_\nu + \frac{\Lambda}{3} \gamma_{\nu\mu} J^\mu. \end{aligned} \quad (24)$$

The terms with Λ are similar to the terms with the Λ in the Lichnerowicz operator ¹ $\Delta_L^{(2)}$ acting on $T^{\mu\nu}$

$$-\Delta_L^{(2)} T^{\mu\nu} = \square T^{\mu\nu} + \frac{2}{3} \Lambda g_{\mu\nu} T - \frac{8}{3} T^{\mu\nu}. \quad (25)$$

One can decompose $\Delta_L^{(2)}$ into a traceless part and a trace, and invert each part separately, just as one decomposes the propagator for gauge fields into a transversal part and a longitudinal part. (One can then prove that the latter does not renormalize). Similarly, one could decompose the spin 3/2 propagator [13].

We now put these results into the context of supersymmetry. For $N = 1$ supergravity, one can add a cosmological constant and still preserve local supersymmetry [9]. The action and transformation rules read

$$\begin{aligned} \mathcal{L} &= \mathcal{L}_E + \mathcal{L}_{RS} + e \left(\frac{6\alpha}{\kappa^4} + \frac{\alpha}{\kappa} \bar{\Psi}_\mu \gamma^{\mu\nu} \Psi_\nu \right) \\ \delta \Psi_\mu &= \frac{1}{\kappa} \nabla_\mu \epsilon + \frac{\alpha}{\kappa^2} \gamma_\mu \epsilon; \quad \delta e_\mu^m = \frac{\kappa}{2} \bar{\epsilon} \gamma^m \psi_\mu, \end{aligned} \quad (26)$$

with α a free constant. We note that the apparent mass term $\frac{\alpha}{\kappa} \bar{\Psi}_\mu \gamma^{\mu\nu} \Psi_\nu$ has the same form as the explicit mass term in (3), but for spin 2 fields the expansion of the cosmological term leads to a different quadratic mass term than the Fierz-Pauli mass term

$$\begin{aligned} \frac{1}{\kappa^4} \sqrt{-\det(g+h)} &= \frac{\sqrt{-g}}{\kappa^4} \left[1 + \frac{1}{2} h - \frac{1}{4} \left(h_{\mu\nu} h^{\mu\nu} - \frac{1}{2} h^2 \right) + \dots \right], \\ \mathcal{L}_{(mass)}^{(2)} &= -\frac{\sqrt{-g} M^2}{8} \left(h_{\mu\nu} h^{\mu\nu} - h^2 \right). \end{aligned} \quad (27)$$

where all the contractions are performed with the background metric $g_{\mu\nu}$. In flat space the massive free spin 2 and spin 3/2 system with Fierz-Pauli mass terms

$$\begin{aligned} \mathcal{L}_{FP} &= -\frac{1}{2} (\partial_\lambda h_{\mu\nu})^2 + (\partial^\nu h_{\mu\nu})^2 - \partial^\mu h \partial^\nu h_{\mu\nu} + \frac{1}{2} (\partial^\mu h)^2 - \frac{1}{2} \bar{\Psi}_\mu \gamma^{\mu\rho\sigma} \partial_\rho \Psi_\sigma \\ &- \frac{M^2}{2} (h_{\mu\nu}^2 - h^2) + \frac{M}{2} \bar{\Psi}_\mu \gamma^{\mu\nu} \Psi_\nu \end{aligned} \quad (28)$$

is not invariant under the rigid susy transformation rules for anti-de Sitter supergravity

$$\delta h_{\mu\nu} = \frac{1}{2} (\bar{\epsilon} \gamma_\mu \Psi_\nu + \bar{\epsilon} \gamma_\nu \Psi_\mu), \quad \delta \Psi_\mu = \frac{1}{4} \hat{\omega}_\mu^{mn} \gamma_{mn} + \alpha M h_{\mu\nu} \gamma^\nu \epsilon + \beta M h \gamma_\mu \epsilon \quad (29)$$

where $\hat{\omega}_{\mu mn} = \frac{1}{8} (-\partial_m h_{\mu n} + \partial_n h_{\mu m})$ is the linearized spin connection.

For the massless case, we can promote the rigid susy to a local susy by using the Noether method and fusing the local symmetry $\delta \Psi_\mu = \partial_\mu \eta$ of the linearized spin 3/2 action with the linearized rigid susy transformation rules. The result is $N = 1$ supergravity [8]. However, to repeat this procedure for the massive theory one runs into a wall. This is not surprising: a massive spin 2 has 5 degrees of freedom and a massive spin 3/2 has 4 degrees of freedom. The massive representation of $N = 1$ susy with spin 2 contains one massive graviton, two gravitinos (one complex gravitino) and one massive real vector field. The degrees of freedom now match:

¹To check the constants, note that if $T_{\mu\nu} = g_{\mu\nu} A$, one has $-\Delta_L^{(2)}(g_{\mu\nu} A) = \square A$.

$5 + 3 = 4 + 4$, and one can begin with a linearized rigid susy system, and use the Noether method to construct the non-linear theory.

One must then not only add a cubic term to the action and quadratic terms to the transformation rules as usual in the Noether method, but also a term linear in h to the action, and a field-independent local term $\delta\Psi_\mu = \gamma_\mu\epsilon$ to the transformation rules. The result is $N = 2$ supergravity with a super-cosmological term [10].

As a final remark, we discuss how to generate masses for the spin 2 Fierz-Pauli action (28) and for the spin 3/2 Rarita-Schwinger action (3) by the Stückelberg formalism [11, 12]. As is well-known, in the case of abelian gauge theories, the Stückelberg formalism provides a gauge invariant procedure to describe a massive vector field. Due to the absence of couplings with the Faddeev-Popov ghosts the theory is renormalizable and the limit $M \rightarrow 0$ is smooth to all orders. The same technique can be used for the linearized Fierz-Pauli action (28) by introducing an auxiliary spin 1 field and an auxiliary spin zero field [3]. It turns out that in flat space the action of the gauge field and of the scalar field are the Maxwell action and the Klein-Gordon action if one chooses the mass terms appropriately. This implies that the action with the auxiliary fields is ghost free.

For the spin 3/2 case, we introduce an auxiliary Majorana spin 1/2 field λ which transforms in the following way

$$\delta\Psi_\mu = \partial_\mu\eta, \quad \delta\lambda = \eta. \quad (30)$$

Therefore, the combination $\Psi_\mu - \partial_\mu\lambda$ is gauge invariant. The spin 3/2 mass term can be written in a gauge invariant way as

$$\mathcal{L}_{mass}^{(3/2)} = \frac{M}{2} (\bar{\Psi}_\mu - \partial_\mu\bar{\lambda}) \gamma^{\mu\nu} (\Psi_\nu - \partial_\nu\lambda) = \frac{M}{2} (\bar{\Psi}_\mu \gamma^{\mu\nu} \Psi_\nu - 2\bar{\Psi}_\mu \gamma^{\mu\nu} \partial_\nu\lambda). \quad (31)$$

The absence of ghosts in the action is due to the cancellation of higher derivatives for the spin 1/2 fields: due to the antisymmetry of the tensor structure $\gamma^{\mu\nu}$ the double derivative term cancels. From this point of view, it becomes clear why a mass term $\bar{\Psi}_\mu\Psi^\mu$ is not allowed: it would lead to a higher-derivative action $\partial_\mu\bar{\lambda}\partial^\mu\lambda$ for the spin 1/2 field. A suitable choice of the gauge fixing is needed in order to remove the couplings $\bar{\Psi}_\mu\gamma^{\mu\nu}\partial_\nu\lambda$, and this generates the Dirac action for the spin 1/2 field. Further details will be discussed in [13].

The structure of the mass terms can also be understood from Kaluza-Klein compactifications from 5 to 4 dimensions. Taking in $\bar{\Psi}_\mu\gamma^{\mu\rho\sigma}\partial_\rho\Psi_\sigma$ the index ρ to be 5, and making the ansatz $\Psi_\sigma(x, x_5) \sim \sqrt{\gamma^5}\psi_\sigma(x)e^{iMx_5}$, one obtains the mass term $M\bar{\psi}_\mu\gamma^{\mu\nu}\psi_\nu$. In a similar manner, one may obtain the Fierz-Pauli mass term in (1) from the Fierz-Pauli action (28) in 5 dimensions by setting $h_{\mu\nu}(x, x_5) = h_{\mu\nu}(x)e^{iMx_5}$ for $\mu, \nu = 0, 3$.

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